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Published in:
Research in Optical Sciences

Publication date:
2012

Document Version
Publisher's PDF, also known as Version of record

[Link back to DTU Orbit](#)

Citation (APA):
Bache, M., & Zhou, B. (2012). Cascaded Soliton Compression of Energetic Femtosecond Pulses at 1030 nm. In *Research in Optical Sciences* (pp. JT2A.60). Optical Society of America.

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Cascaded soliton compression of energetic femtosecond pulses at 1030 nm

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Abstract: We discuss soliton compression with cascaded second-harmonic generation of energetic femtosecond pulses at 1030 nm. We discuss problems encountered with soliton compression of long pulses and show that sub-10 fs compressed pulses can be achieved.

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OCIS codes: (320.5520); (320.7110); (190.5530); (190.2620); (320.2250).

Ultrafast fiber laser systems using Yb-doped laser materials are rugged, cheap and compact platforms that currently are undergoing a rapid development, and by employing the chirped pulse amplification technique high-energy femtosecond pulses can be generated with μJ –sub-mJ pulse energies. In the same way solid-state Yb-based amplifiers can give multi-mJ pulse energies and interesting alternatives to Ti:Sapphire systems because they can be diode-pumped. However, the limited gain bandwidth of Yb systems implies an amplified pulse duration of around 500 fs, and efficient post-compression methods are therefore needed to reach sub-100 fs. The standard way of compressing energetic pulses today is to use 0.5-1.0 m long hollow gas-filled fibers [1], where spectral broadening is achieved through the formation of a filament using the self-focusing nonlinearities of the gas. Temporal compression is then achieved by a pair of gratings. With this technique compression from 500 fs to 60 fs was observed at $\lambda = 1.03 \mu\text{m}$ [2].

We here propose to use cascaded (phase-mismatched) second-harmonic generation as pulse compressor. With this technique high-energy pulse compression of the pump pulse towards few-cycle duration has been achieved [3–7]. These studies were motivated by the possibility of creating a negative (self-defocusing) Kerr-like nonlinearity. Thus, self-focusing problems encountered in Kerr-based compressors are avoided, and the input pulse energy is practically unlimited [3]. Moreover, small-scale self-focusing does not occur either since there is no spatial modulation instability gain for self-defocusing nonlinearities. Furthermore both spectral broadening and temporal compression can occur in a single nonlinear material through solitons [4]. Solitons are stable nonlinear waves that exist as a balance between nonlinearity and dispersion, which for a negative nonlinearity is achieved with normal dispersion. Thus, in cascaded SHG solitons can form in the near-IR where the majority of lasers operate and where most materials have normal dispersion. Here we investigate cascaded soliton compression at the operating wavelength of Yb-based laser amplifiers in standard nonlinear crystals (β -barium-borate, BBO and lithium niobate, LN). We first show that compression of 500 fs pulses towards few-cycle duration is possible in BBO, but that it requires very long interaction lengths so spatial walk-off becomes an issue. Assuming as input 60 fs compressed pulses from a hollow-fiber compressor, we show pulse compression to few-cycle duration in few-mm short LN crystals, where spatial walk-off is absent.

In cascaded SHG the conversion process L is strongly phase-mismatched $|\Delta k L| \gg 1$: pump up-conversion to the second harmonic (SH) is after a coherence length $\pi/|\Delta k|$ followed by the reverse process of down-conversion to the pump. On continued propagation the SH is therefore cyclically generated and back-converted. In this cascaded nonlinear interaction the pump experiences a nonlinear phase shift due to the difference in phase velocities, and the magnitude and sign of the phase shift is determined by the phase-mismatch parameter $\Delta k = k_2 - 2k_1$, where $k_j = k_j(\omega_j)$, $k_j(\omega) = n_j(\omega)\omega/c$ is the wavenumber, and $n_j(\omega)$ is the linear refractive index. To see this, we note that in the cascading limit the coupled SHG equations can approximately be reduced to a single equation for the pump. Neglecting for simplicity higher-order dispersion, self-steepening effects and Kerr cross-phase modulation (XPM), and assuming normal group-velocity dispersion (GVD) $k_1^{(2)} \equiv d^2 k_1(\omega)/d\omega^2|_{\omega=\omega_1} > 0$ it is [8–10]

$$\left(i\frac{\partial}{\partial \xi} - \frac{\partial^2}{\partial \tau^2}\right) U_1 + N_{\text{Kerr}}^2 \left(U_1 |U_1|^2 - \tau_R U_1 \frac{\partial |U_1|^2}{\partial \tau}\right) - s_{\Delta k} N_{\text{casc}}^2 \left(U_1 |U_1|^2 + i \text{sgn}(d_{12}) \tau_{R,\text{casc}} |U_1|^2 \frac{\partial U_1}{\partial \tau}\right) = 0 \quad (1)$$

where $s_{\Delta k} = \text{sgn}(\Delta k)$. This approximate equation includes material Kerr self-phase modulation (SPM) as well as a delayed Raman nonlinearity (here approximated for a long pulse, i.e. narrow spectrum) with the normalized Raman time $\tau_R = T_R/T_{\text{in}}$. For most materials $T_R < 5$ fs. The normalization is chosen to give soliton units, so $\xi = z/L_{D,1}$,

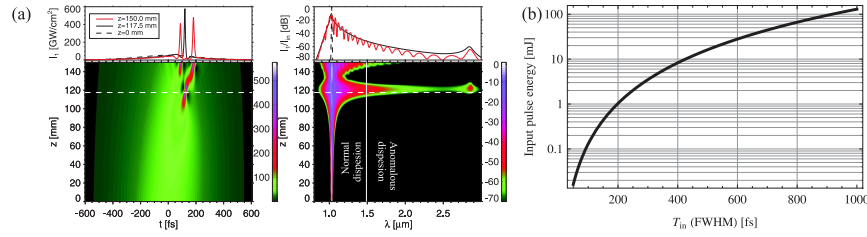


Figure 1. (a) Simulation of soliton pulse compression in BBO, showing the propagation dynamics of the pump vs. time (left) and wavelength (right). The input pulse is a $\lambda_1 = 1.03 \mu\text{m}$ sech-shaped 500 fs FWHM pulse and $I_{\text{in}} = 54 \text{ GW/cm}^2$, giving $N_{\text{eff}} = 10$, $N_{\text{casc}} = 20.3$ and $N_{\text{Kerr}} = 17.7$. The simulation used the full coupled SHG slowly-evolving envelope equations in the plane-wave limit [8] including self-steepening, Kerr SPM and XPM effects, exact chromatic dispersion using the Sellmeier equations, and using $n_{\text{Kerr}}^I = 5.8 \times 10^{-20} \text{ m}^2/\text{W}$, $\theta = 20.2^\circ$, $\Delta k = 60 \text{ mm}^{-1}$ and $d_{\text{eff}} = 2.09 \text{ pm/V}$, $f_R = 0$ (see [14] for more details). (b) The minimum input pulse energy required to avoid spatial walk-off during propagation, shown vs. the input pulse FWHM.

$\tau = t/T_{\text{in}}$ and $|U_1|^2 = I_1/I_{\text{in}}$. This gives a Kerr soliton order $N_{\text{Kerr}}^2 = L_{D,1} \frac{\omega_1}{c} n_{\text{Kerr}}^I I_{\text{in}}$, where $L_{D,1} = T_{\text{in}}^2/|k_1^{(2)}|$ is the pump dispersion length, T_{in} is the input pulse duration, $n_{\text{Kerr}}^I > 0$ is the self-focusing material Kerr nonlinear refractive index, and I_{in} is the peak input intensity. The cascaded nonlinearity is shown under the weakly nonlocal approximation [9], which holds when the phase mismatch is large enough to be in the *stationary regime* $\Delta k > \Delta k_{\text{sr}} = d_{12}^2/2k_2^{(2)}$. The cascading gives a SPM term whose strength is given by the cascading soliton order $N_{\text{casc}}^2 = L_{D,1} \frac{\omega_1}{c} |n_{\text{casc}}^I| I_{\text{in}}$, where the cascading nonlinear refractive index is $n_{\text{casc}}^I \simeq -4\pi d_{\text{eff}}^2/c\epsilon_0\lambda_1 n_1^2 n_2 \Delta k$ [11], and d_{eff} is the effective quadratic nonlinearity. Note that the strength and sign is controlled by Δk , and for $\Delta k > 0$ the cascading term is self-defocusing. The cascading also gives a Raman-like perturbation of the pump, which is characterized by the normalized parameter $\tau_{R,\text{casc}} = T_{R,\text{casc}}/T_{\text{in}} = 2|d_{12}|/T_{\text{in}}\Delta k$ [9, 12]. This effect stems from group-velocity mismatch (GVM), $d_{12} = k_1^{(1)} - k_2^{(1)}$ is the GVM parameter, and results in a steepening-like effect [13]. Finally, Eq. (1) shows that the effective SPM can be expressed by an *effective* nonlinear refractive index $n_{\text{eff}}^I = n_{\text{casc}}^I + n_{\text{Kerr}}^I$. For the case $\Delta k > 0$ and $n_{\text{eff}}^I < 0$ we can introduce an effective soliton order $N_{\text{eff}}^2 = N_{\text{casc}}^2 - N_{\text{Kerr}}^2$, which characterizes the cascaded soliton behavior [8].

We have previously studied cascaded soliton pulse compression at $\lambda_1 \simeq 1.03 - 1.06 \mu\text{m}$: in Ref. [8, 10] we used a critical (type-I) phase-mismatched BBO nonlinear crystal through an $oo \rightarrow e$ interaction. We show in Fig. 1(a) a BBO simulation using a 500 fs FWHM sech-shaped pulse as input. The phase mismatch $\Delta k = 60 \text{ mm}^{-1}$ was chosen so compression occurs in the stationary regime ($\Delta k_{\text{sr}} = 48 \text{ mm}^{-1}$), and $n_{\text{casc}}^I = -7.4 \times 10^{-20} \text{ m}^2/\text{W}$. The chosen soliton order $N_{\text{eff}} = 10$ should give a compression factor $f_c = 4.7(N_{\text{eff}} - 0.86) = 43$ [8], and indeed the pulse compresses to 11 fs FWHM (3 optical cycles) after 120 mm. At the compression point a dispersive wave forms in the anomalous (linear) dispersion regime [14], and after this pulse splitting is observed. Thus, strong pulse compression to few-cycle duration is possible in BBO at $\lambda_1 = 1.03$, but the long compression length is problematic since it is much longer than standard crystal sizes. The problem can be explained from observing the optimal compression point z_{opt} , i.e. the point where the soliton pulse duration is the shortest. For large effective soliton orders it scales as [8]

$$z_{\text{opt}} \simeq L_{D,1} \frac{\pi}{2} 0.44/N_{\text{eff}} = 0.44 T_{\text{in}} \sqrt{\pi \lambda_1 / [8k_1^{(2)} |n_{\text{eff}}^I| I_{\text{in}}]}, \quad N_{\text{eff}} \gg 1 \quad (2)$$

Thus long pulses, low GVD and a low effective nonlinearity all result in an increased compression length. It can be reduced by taking a larger N_{eff} (i.e. intensity), but the price is a dramatic increase in the GVM-induced steepening ($T_{R,\text{casc}} = 3.8 \text{ fs}$ here, and to increase N_{eff} one must increase N_{casc} even more). Secondly, the spatial walk-off angle of BBO (for the present case $\rho = 2.8^\circ$) was obviously not modeled in the plane-wave simulation. The walk-off length is $L_{\text{wo}} = w_0/\tan\rho$, where w_0 is the Gaussian waist of the beam. We remind that for a sech-shaped pulse, the pulse energy is $W_{\text{in}} = \pi w_0^2 T_{\text{in}} I_{\text{in}}$. Consider a 2 mJ pulse energy: the chosen intensity of Fig. 2(a) is achieved with $w_0 = 2.0 \text{ mm}$, and thus $L_{\text{wo}} = 40 \text{ mm} \ll z_{\text{opt}}$. With a higher pulse energy we could increase w_0 and thereby L_{wo} , but keep the same intensity. Let us estimate how much pulse energy is needed to avoid walk-off effects by requiring $z_{\text{opt}} = L_{\text{wo}}$ and use Eq. (2) to get $W_{\text{in}} (\text{min}) = \tan^2 \rho 0.44^2 \pi^2 \lambda_1 T_{\text{in}}^3 / [8k_1^{(2)} |n_{\text{eff}}^I|]$, where for a sech-shaped pulse $T_{\text{in}} = T_{\text{in}}^{\text{FWHM}}/2\ln(1+\sqrt{2})$. This is shown in Fig. 1(b) for the present case. A 500 fs FWHM pulse requires 16 mJ input energy while much less energy is needed for shorter input pulses. We finally mention that both problems could be solved by having several crystals in a row (increased length) and flipping the optical axis for each segment to compensate for the walk-off [15].

We could also choose a crystal with much higher GVD, such as LN, where GVD is app. 5 times larger than BBO. In Ref. [16] we investigated LN in a type-I interaction, but unfortunately compression in the stationary regime was not possible because achieving $n_{\text{eff}}^I < 0$ required that $\Delta k < \Delta k_{\text{sr}}$. We solved this problem recently by using LN in a

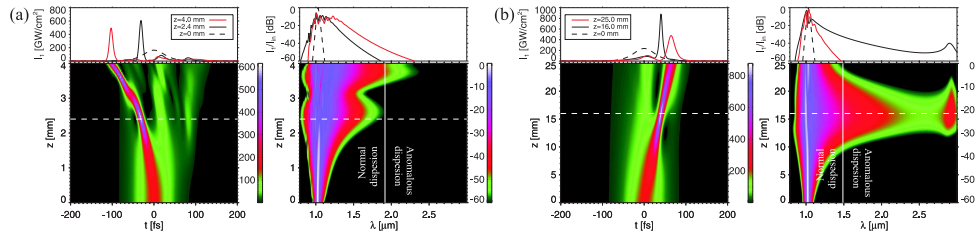


Figure 2. Cascaded soliton compression with $T_{in} = 60$ fs FWHM and $N_{eff} = 2.5$ in (a) type-0 LN (b) type-I BBO [as in Fig. 1(a) except for $I_{in} = 235$ GW/cm², $N_{casc} = 5.1$ and $N_{Kerr} = 4.4$]. LN simulation: $I_{in} = 160$ GW/cm², $N_{casc} = 3.7$ and $N_{Kerr} = 3.9$, $\theta = 90^\circ$, $d_{eff} = 25$ pm/V, $\Delta k = 998$ mm⁻¹, $n_{Kerr}^I = 30 \times 10^{-20}$ m²/W, and $f_R = 0.51$ (see [7] for more details).

noncritical (type 0, $ee \rightarrow e$) phase-matching configuration, which has such a large d_{eff} that $n_{eff}^I < 0$ despite a huge phase mismatch. At the same time the large Δk exactly ensured $\Delta k > \Delta k_{sr}$. We showed compression of 50 fs FWHM pulses with $\lambda_1 = 1.3$ μ m to 16 fs [7], and the huge GVD in LN meant that compression occurred after just 1 mm. In such noncritical interaction the spatial walk-off is also absent. We estimated from simulations that $n_{Kerr}^I = 30 \times 10^{-20}$ m²/W. Unfortunately for $\lambda_1 = 1.03$ μ m, even if the cascading is stationary $\Delta k > \Delta k_{sr} = 490$ mm⁻¹, we can only get $n_{casc}^I = -28 \times 10^{-20}$ m²/W. Thus, $n_{eff}^I > 0$ and no solitons should form, which we have verified numerically. However, LN has a quite strong Raman response: we estimated in Ref. [7] $f_R = 0.51$. For a shorter input pulse this means that the Kerr term in the second bracket in Eq. (1) becomes inaccurate, and must be replaced by the more accurate term $(1 - f_R)U_1|U_1|^2 + f_R U_1 \int_{-\infty}^{\infty} dt' R(t')|U_1(\tau - t')|^2$. This states that the Kerr SPM term experienced by a short pulse is $1 - f_R$ smaller than for a long pulse. Since $n_{eff}^I = n_{casc}^I + (1 - f_R)n_{Kerr}^I < 0$ in LN at $\lambda_1 = 1.03$ μ m, solitons can now form using short input pulses. Fig. 2(a) shows a full simulation with $N_{eff} = 2.5$ and a 60 fs FWHM input pulse (viz a 500 fs laser pulse compressed in a hollow-fiber [2]). Indeed, clear soliton compression is observed after only 2.5 mm to 10 fs. Some pulse splitting is seen, which is due to Raman effects [7] that also cause the spectrum to red shift. Compare this with (b) showing a similar compression in BBO: the same pulse duration is achieved, but after 5 times longer propagation. Instead the absence of Raman effects in the BBO simulation gives a cleaner pulse.

Acknowledgments: the Danish Council for Independent Research (*Femto-VINIR* project, 274-08-0479)

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